Melting or nucleon transfer in fusion of heavy nuclei ?

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Abstract

The time-dependent transition between a diabatic interaction potential in the entrance channel and an adiabatic potential during the fusion process is investigated within the two-center shell model. A large hindrance is obtained for the motion to smaller elongations of near symmetric dinuclear systems. The comparison of the calculated energy thresholds for the complete fusion in different relevant collective variables shows that the dinuclear system prefers to evolve in the mass asymmetry coordinate by nucleon transfer to the compound nucleus.

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The experimental synthesis of new superheavy elements [1,2] and nuclei far from the line of stability stimulates the study of fusion processes in heavy ion collisions at low energies (< 15 MeV/nucleon). The experimental results on the fusion of heavy nuclei and the cold fusion reactions can not be satisfactorily explained as a melting of two nuclei in the relative coordinate with the known macroscopic and microscopic models [3–5]. In these models the fusion cross section is overestimated for near symmetric reactions, for example, for the ^{110}Pd+^{110}Pd reaction, and the optimal bombarding energy and isotopic dependence of the fusion probabilities are incorrect [1,2]. In the dinuclear system (DNS) concept [6,7] the fusion is explained as a transfer of nucleons or clusters between nuclei in a dinuclear (quasimolecular) configuration. Therefore, fusion occurs in the degree of freedom of mass asymmetry $\eta = (A_2 - A_1)/A$ (A_1 and A_2 are the mass numbers on both sides of the neck and $A = A_1 + A_2$), and this model describes the fusion in symmetric reactions with heavy nuclei and in reactions producing superheavy nuclei quite well. These facts stimulate the microscopical justification of relevant collective coordinates which are mainly responsible for the complete fusion. It will be interesting to study the competition between two possible fusion channels. The first one (λ–channel) describes the transition of two nuclei into the compound nucleus with the elongation coordinate λ of the two–center shell model (TCSM) which measures the length l of the two–nucleus system in units of the diameter $2R_0$ of a spherical assumed compound nucleus: $\lambda = l/2R_0$. This channel assumes a fixed mass asymmetry during the fusion and is named λ–channel. The second channel, named η–channel, describes the evolution of the DNS to the compound nucleus as a change of the mass asymmetry η by nucleon transfer from the light nucleus to the heavy one (the DNS concept [6,7]). The comparison of the fusion probability calculated in both channels will allow us to find the favorable fusion channel.

In microscopic studies of the dynamics of fusion, based on group theory and on the TCSM, we obtained a large structural forbiddenness for the fast growth of the neck and for the motion to smaller elongations in the dinuclear system [8–10]. The effect of structural forbiddenness is the reason for the stability of the DNS configurations against a dissolution into the λ–channel. The decrease of the structural forbiddenness with increasing mass asymmetry [9] supports the idea of the DNS concept [6,7] that complete fusion takes place after the mass asymmetry η is increased by thermal fluctuations. In the present letter we will study whether the system has time for destroying the "memory" on the structural forbiddenness. This time is necessary to reorganize the density of the system for the transition from the initial diabatic potential $V_{di}(\lambda)$ to the adiabatic potential $V_{ad}(\lambda)$. The value of $V_{di}(\lambda)$ could be represented by a frozen density approximation, but the diabatic potential and the frozen density potential are conceptually and physically not equivalent even though the numerical values are similar [11,12]. The repulsive character of the potential is mainly related to the diabatic particle–hole excitations and to the compression effects in diabatic and frozen density considerations, respectively [11]. The time dependence of the transition of the potential can be related to the effective characteristic relaxation time τ for the shape degrees of freedom of the system [13,14]

$$V(\lambda, t) = V_{di}(\lambda) \exp(-\int_0^t \frac{dt}{\tau(\lambda, t)}) + V_{ad}(\lambda)[1 - \exp(-\int_0^t \frac{dt}{\tau(\lambda, t)})]. \qquad (1)$$

A time-dependent dynamical potential $V(\lambda, t)$ was originally introduced in Refs. [13,14] from

a phenomenological ansatz and applied to study the effects of local equilibrium in dissipative heavy-ion collisions. The Eq. (1) may be rewritten as

$$V(\lambda, t) = V_{ad}(\lambda) + \Delta V_{di}(\lambda, t) \qquad (2)$$

with $\Delta V_{di}(\lambda, t) = (V_{di}(\lambda) - V_{ad}(\lambda)) \exp(-\int_0^t \frac{dt}{\tau(\lambda, t)})$. The additional part $\Delta V_{di}(\lambda, t)$ can be microscopically obtained from the diabatic excitation of particle-hole states

$$\Delta V_{di}(\lambda, t) \approx \sum_\alpha \epsilon_\alpha^{di}(\lambda)[n_\alpha^{di}(\lambda, t) - n_\alpha^{ad}(\lambda)], \qquad (3)$$

where the $\epsilon_\alpha^{di}(\lambda)$ are diabatic single–particle energies as a function of the elongation λ of the TCSM. The adiabatic occupation numbers $n_\alpha^{ad}(\lambda)$ vary with λ according to a Fermi distribution with the temperature $T(\lambda) = \sqrt{E^*(\lambda)/a}$ ($a = A/12$ MeV^{-1}), where the excitation energy $E^*(\lambda)$ is determined from total energy conservation. The exponential factor in (1) is due to the dependence of the diabatic occupation probabilities n_α^{di} on time expressed by the relaxation equations [15–17]

$$\frac{dn_\alpha^{di}(\lambda, t)}{dt} = -\frac{1}{\tau(\lambda, t)}[n_\alpha^{di}(\lambda, t) - n_\alpha^{ad}(\lambda)], \qquad (4)$$

which is known in the relaxation time approximation. Due to the residual two-body interactions, the diabatic occupation probabilities approach a local (fixed λ) equilibrium with an average relaxation time

$$\tau(\lambda, t) = \frac{2\hbar}{< \Gamma(\lambda, t) >}. \qquad (5)$$

The factor 2 assumes that two subsequent collisions are sufficient to establish equilibrium for fixed values of the collective variable λ. Here, we use a minimal value of this factor (or minimal possible value of τ) in comparison to Refs. [15–17] where this factor was chosen as 3–4. From the Eq. (1) it is clear that the effective time τ necessary to reorganized the densities of the system corresponds to a mean value of various relaxation times associated to the shape degrees of freedom of system which is larger than average single-particle decay time $(\frac{\hbar}{<\Gamma>})$ due to the effect of self–consistency between collective and single-particle degrees of freedom [18,19]. The width in Eq. (5)

$$< \Gamma(\lambda, t) >= \sum_\alpha \overline{n}_\alpha^{di}(\lambda, t)\Gamma_\alpha(\lambda)/\sum_\alpha \overline{n}_\alpha^{di}(\lambda, t) \qquad (6)$$

is an average width of the particle-states above the Fermi level ($\overline{n}_\alpha^{di} = n_\alpha^{di}$ for $\epsilon_\alpha^{di} > \epsilon_F$) and of the hole-states under the Fermi level ($\overline{n}_\alpha^{di} = 1 - n_\alpha^{di}$ for $\epsilon_\alpha^{di} \leq \epsilon_F$). For the widths Γ_α, the following expression is used

$$\Gamma_\alpha = \Gamma_0^{-1}\frac{(\epsilon_\alpha^{di} - \epsilon_F)^2}{1 + [(\epsilon_\alpha^{di} - \epsilon_F)^2]/c^2}. \qquad (7)$$

The parameters Γ_0 and c are known from studies with the optical model potential and the effective masses, and their values are in the range 0.030 MeV$^{-1} \leq \Gamma_0^{-1} \leq 0.061$ MeV^{-1} and

15 MeV$\leq c \leq 30$ MeV. The results depend weakly on the value of the parameter c. In the calculations we take the standard value c=20 MeV and consider the cases with the two extreme values of Γ_0^{-1} mentioned above. From an extended expression of Eq. (7) one can see for very large free energies $\epsilon_\alpha^{di} - \epsilon_F$ that the broadening of single-particle widths due to intrinsic excitation energy of system plays no essential role in contrast to the case when the excited system is near the equilibrium state [18,20,21]. Although one may define a local excitation energy during the decay of the diabatic potential to the adiabatic one, the concept of temperature is less meaningful as the system is not locally equilibrated or thermalized [22]. More detailed investigations are required to clarify these points.

Since in fusion and quasifission we deal with two-center systems, we use the two-center shell model (TCSM) [23–25] for calculating the adiabatic or diabatic potential energy surface. In the TCSM the nuclear shapes are defined by a set of coordinates. The elongation λ measures the length l of the system and is used to describe the relative motion. The transition of the nucleons through the neck is described by the mass asymmetry η. The neck parameter $\varepsilon = E_0/E'$ is defined by the ratio of the actual barrier height E_0 to the barrier height E' of the two-center oscillator. The deformations $\beta_i = a_i/b_i$ of axial symmetric fragments are related to the ratio of their semiaxes.

With the TCSM and the Strutinsky method the adiabatic potential energy V_{ad} can be calculated as the sum of a macroscopic energy U_{LDM} obtained by the liquid drop model, microscopic correction δU_{shell} that arises due to the shell structure of the nuclear system, energy of the pairing correction δU_{pair} and the proximity nuclear potential V_N to improve the values of the adiabatic energy obtained for large elongations [10,12]

$$V_{ad} = U_{LDM}(\lambda, \varepsilon, \eta) + V_N(\lambda, \varepsilon, \eta) + \delta U_{shell}(\lambda, \varepsilon, \eta, E^*(\lambda)) + \delta U_{pair}(\lambda, \varepsilon, \eta, E^*(\lambda)). \quad (8)$$

We neglect the dependence of potential energy on the angular momentum in the reactions considered [7] because in fusion reactions with heavy nuclei only low angular momenta ($<20-30\,\hbar$) contribute [1,2,26]. Shell effects are damped exponentially $\delta U_{shell} = \delta U_{shell}(E^* = 0)\exp(-\zeta E^*)$ if the system is excited with the excitation energy E^*. The parameter ζ is chosen as $\zeta^{-1} = 5.48A^{\frac{1}{3}}/(1 + 1.3A^{-\frac{1}{3}})$ MeV [27]. The pairing corrections are taken as follows δU_{pair}=0 for $E^* \geq E_c$ and $\delta U_{pair}=\delta U_{pair}(E^* = 0)[1 - E^*/E_c]^2$ for $E^* < E_c$ [2] with $E_c = 10 MeV$. The isotopic composition of the nuclei forming the DNS was chosen with the condition of the N/Z-equilibrium in the system [7]. For the touching configuration of nuclei the diabatic effects are very small and the diabatic potential practically coincides with the adiabatic one [12]. So, there is no any difference between the DNS evolution in η in the adiabatic and diabatic basises. Because of this the values of fusion barrier in η–channel and quasifission barrier are practically independent of time. For the smaller elongations $\lambda < \lambda_t$ the diabatic potential is considerably larger than the adiabatic potential.

The diabatic levels ϵ_α^{di} are classified by the quantum numbers $\alpha = j_z, l_z, s_z, n_\rho, n_z$ of the eigenstates of the diabatic Hamiltonian $H' = H - SV$ which is obtained by the method of maximum symmetry [16,12]. The method of maximum symmetry eliminates the symmetry violating parts SV from the total Hamiltonian H of the TCSM. Diabatic levels obtained by this method agree with those of the maximum overlap procedure [16]. We use the method of maximum symmetry because it is numerically easy to handle. In most reactions studied in [12] the diabatic potential has a minimum in the neck parameter ε around $\varepsilon = 0.65$– 0.85. Nuclei are considered as spherical with $\varepsilon = 0.74$ which corresponds to realistic shapes

of the DNS for λ =1.5–1.6. Due to the large inertia and friction coefficients in the neck coordinate ε obtained in a microscopical treatment [10], the DNS configurations with fixed neck parameters have a long lifetime in comparison to the reaction time.

In order to study the competition between the fusion channels in λ and η, we use the fusion rate $\Lambda_{fus}^{\lambda}(t)$ $(\Lambda_{fus}^{\eta}(t))$ through the inner fusion barrier B_{fus}^{λ} (B_{fus}^{η}) in λ (η) to calculate the fusion probability in λ–channel (η–channel)

$$P_{fus}^{\lambda(\eta)} = \int\limits_{0}^{t_0} \Lambda_{fus}^{\lambda(\eta)}(t)dt, \tag{9}$$

where the lifetime t_0 of the DNS is obtained with the condition

$$\int\limits_{0}^{t_0} [\Lambda_{fus}^{\lambda}(t) + \Lambda_{fus}^{\eta}(t) + \Lambda_{qf}^{\lambda}(t)]dt = 1. \tag{10}$$

The rate of probability $\Lambda_{qf}^{\lambda}(t)$ through the external barrier in λ determines the quasifission process (the decay of the system). The height B_{qf}^{λ} of this barrier monotonically decreases with the DNS mass asymmetry (at fixed total mass and charge numbers of system) because the increase of the Coulomb repulsion with decreasing η leads to very shallow pocket in nucleus–nucleus potential. So, the quasifission probability for the symmetric and near symmetric DNS is much larger than for asymmetric one. Because the reactions considered in paper are symmetrical or near symmetrical and, correspondingly, the initial DNS of them are in or near the minimum of the total potential energy of system as function of λ and η, the main contribution to quasifission channel comes from the initial or near initial configurations which have more or less the same quasifission barrier B_{qf}^{λ}. All these facts allow us to calculate the quasifission rate for the initial DNS using the Kramers-type expression. This decay process in λ determines mainly the lifetime of the DNS, because the barrier B_{qf}^{λ} is smaller than the barrier B_{fus}^{η} in mass asymmetry. The lifetimes t_0 obtained for the reactions considered are comparable with the experimentally extracted characteristic fusion times of $10^{-21} - 10^{-20} s$ [28] . Since the rates assume their final values very fast [7,10] and their initial contributions are of the order of the accuracy of the calculation of the barrier heights, we use the one-dimensional Kramers expression [29] ($^{Kr}\Lambda_i^j$, $i =$"fus" or "qf" and $j =$"λ" or "η") which is a quasi–stationary solution of the Fokker-Planck equation for the corresponding rate of probability

$$^{Kr}\Lambda_i^j = \frac{\omega_j}{2\pi\omega^{B_i^j}} \left(\sqrt{(\frac{g}{2\hbar})^2 + (\omega^{B_i^j})^2} - \frac{g}{2\hbar} \right) \exp(-\frac{B_i^j}{T(\lambda_t)}). \tag{11}$$

Here, B_i^j denotes the height of the fusion barriers $B_{fus}^{\lambda}(t)$ and B_{fus}^{η}, and quasifission barrier B_{qf}^{λ}. The values of B_{fus}^{η} and B_{qf}^{λ} are practically independent of time. The initial DNS of the reactions considered are in or near the minimum of the total potential energy of system as function of λ and η as mentioned above and are in thermodynamic equilibrium because diabatic and adiabatic potentials practically coincides. The temperature $T(\lambda_t)$ is calculated by using the expression $T(\lambda_t) = \sqrt{E^*(\lambda_t)/a}$ where λ_t is the elongation of the touching nuclei, i.e. of the initial DNS. In the calculations we assume that the excitation

energy of the initial DNS is $E^*(\lambda_t)$=30 MeV in all reactions considered. In (11), $\omega^{B_i^j}$ is the frequency of the inverted harmonic oscillator approximating the potential in the variable j on the top of the fusion or quasifission barriers B_i^j, and ω_j is the frequency of the harmonic oscillator approximating the potential in the variable j for the initial DNS. The method of the calculation of the mass parameters is given in Refs. [7,10].

In our calculations of the fusion probabilities, the following values are used $\hbar\omega^{B_{qf}^\lambda} \approx$ 0.8–1.0 MeV, $\hbar\omega^{B_{fus}^\eta} \approx$ 1.5–2.0 MeV, $\hbar\omega_\lambda \approx$ 1.5–2.0 MeV and $\hbar\omega_\eta \approx$ 0.8–1.0 MeV for the reactions considered [7]. The value of $\hbar\omega^{B_{fus}^\lambda}$ at the inner fusion barrier in λ is about 0.5–0.6 MeV and agrees with the one obtained in the calculations of fission [30] at the saddle point. The friction coefficients obtained with g=2 MeV have the same order of magnitude as the ones calculated within the one-body dissipation models [7,10]. The values obtained for $P_{fus}^{\lambda(\eta)}$ depend rather weakly on g in (11) [7]. The possibility to apply the Kramers expression to relatively small barriers was demonstrated in [31].

The time-dependent diabatic potentials for the reactions ^{110}Pd+^{110}Pd and ^{124}Sn+^{124}Sn are presented in Figs. 1a) and b). The time-dependent inner fusion barrier B_{fus}^λ in λ appears due to the dependence of the relaxation time of the diabatic potential on the elongation λ as shown in Fig. 2 for the ^{110}Pd+^{110}Pd reaction. The decrease of $<\Gamma>$ in time causes a slower transition of the diabatic potential to the adiabatic potential when this potential approaches the adiabatic limit. The structures in the vanishing diabatic potential (Fig. 1a) are caused by the structures in $<\Gamma>$ as a function of λ (Fig.2) which disappear in time more rapidly than the structures of the diabatic potential. Fig. 3a shows the dependence of B_{fus}^λ on time for the reactions ^{110}Pd+^{110}Pd (η=0) and ^{56}Cr+^{164}Er (η=0.5) which produce the same compound nucleus ^{220}U. The inner fusion barrier in λ for the asymmetric DNS appears later and is pronounced smaller than the one for the symmetric DNS and decreases slower in time. The smaller values of B_{fus}^λ for the asymmetric DNS can be explained by the structural forbiddenness for the motion to smaller values of λ which is smaller than the hindrance in the symmetric case [9,12]. The lifetime t_0 of the DNS formed in both reactions is about $8\cdot10^{-21}$s and the values of B_{fus}^λ at this time are larger than the corresponding fusion barriers B_{fus}^η in η (see Fig. 3b and Table 1). So, the fusion probability P_{fus}^λ in λ is smaller than P_{fus}^η in η (Table 2). This is also demonstrated in Tables 1 and 2 for the reactions ^{123}Sn+^{123}Sn, ^{110}Pd+^{136}Xe, ^{86}Kr+^{160}Gd and ^{76}Ge+^{170}Er which lead to the same compound nucleus ^{246}Fm. The calculated values of P_{fus}^η are in agreement with fusion probabilities extracted from the experimental data [32]. The fusion barrier along mass asymmetry does not depend on time because the diabatic potential energy at the touching configurations with different η is very close to the adiabatic potential energy. From Tables 1 and 2 one can see that the fusion probability P_{fus}^λ increases with increasing mass asymmetry in the entrance channel. It follows from our analysis that in the λ-channel as well as in the η-channel the complete fusion in symmetric reactions yields smaller cross sections in comparison with asymmetric combinations .

As shown in Figs. 1a) and b), the value of $B_{fus}^\lambda(t_0)$ for the reaction ^{124}Sn+^{124}Sn is larger than the value for ^{110}Pd+^{110}Pd. An increasing mass number of the system generally causes an increase of the repulsive character of the initial diabatic potential [12] and decreases the value of the quasifission barrier which mainly determines the DNS lifetime t_0. The same behaviour of the fusion probability P_{fus}^λ is obtained for the symmetric reactions ^{90}Zr + ^{90}Zr, ^{100}Mo + ^{100}Mo, ^{110}Pd + ^{110}Pd, ^{123,124}Sn + ^{123,124}Sn and ^{136}Xe + ^{136}Xe as well (Table 2).

Despite of the strong decrease of B_{fus}^{λ} with the change of the parameter Γ_0 from the maximal to the minimal value, the fusion probabilities obtained in the λ–channel remain to be much smaller than the fusion probabilities obtained in the η–channel which are similar to the experimental values [2,32] (see Table 2). In the heavier system the difference between the fusion barriers and probabilities in both channels is larger and the λ–channel is practically closed. This means a dominance of the fusion in the mass asymmetry degree of freedom which is the fundamental assumption in the DNS concept [6,7].

We discussed which of the two approximations, the diabatic or the adiabatic, is closer to reality in the fusion process. The actual situation depends very much on the initial diabatic potential and on the ratio of the quasifission time or lifetime t_0 of the DNS and the equilibrium time τ. The time–dependent transition between diabatic and adiabatic potentials is a slower process than the quasifission one and the system has not enough time for destroying the "memory" on the structural forbiddenness which is in agreement with estimations in Ref. [9]. As the result, a large hindrance for the motion to smaller elongations λ of the DNS is obtained. The comparison of the calculated energy thresholds for the complete fusion in the λ– and η–channels shows that the DNS favorably evolves to the compound nucleus in mass asymmetry due to the thermal fluctuations.

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TABLE I. Quasifission and inner fusion barriers in η and λ calculated within the TCSM for various symmetric and asymmetric reactions. The inner fusion barriers in λ are given for the lifetimes t_0 of the DNS formed in these reactions. The notations 1) and 2) mean that the values of $B_{fus}^\lambda(t_0)$ are calculated with $\Gamma_0^{-1} = 0.030$ MeV^{-1} and 0.061 MeV^{-1}, respectively.

Reactions	B_{qf}^λ [MeV]	B_{fus}^η [MeV]	t_0 [10^{-21} s]	1) $B_{fus}^\lambda(t_0)$ [MeV]	2) $B_{fus}^\lambda(t_0)$ [MeV]
^{90}Zr+^{90}Zr $\rightarrow$^{180}Hg	2.9	6	20	10	4
^{100}Mo+^{100}Mo $\rightarrow$^{200}Po	2.2	8	15	12	5
^{110}Pd+^{110}Pd $\rightarrow$^{220}U	1.3	12	8	36	14
^{56}Cr+^{164}Er $\rightarrow$^{220}U	2.6	2	8	14	4
^{76}Ge+^{170}Er $\rightarrow$^{246}Fm	0.4	10	5	53	27
^{86}Kr+^{160}Gd $\rightarrow$^{246}Fm	0.2	12	4	65	39
^{110}Pd+^{136}Xe $\rightarrow$^{246}Fm	0.1	15	3	91	54
^{123}Sn+^{123}Sn $\rightarrow$^{246}Fm	0.1	16	3	112	67
^{136}Xe+^{136}Xe $\rightarrow$^{272}Hs	0	22	2	237	154

TABLE II. Fusion probabilities $P_{fus}^{\lambda,\eta}$ in the $\lambda-$ and $\eta-$ channels calculated for the reactions presented in the Table 1 are compared with known experimental values $P_{fus}^{\exp}$ [2,7,10,32]. The notations 1) and 2) are the same as in Table 1.

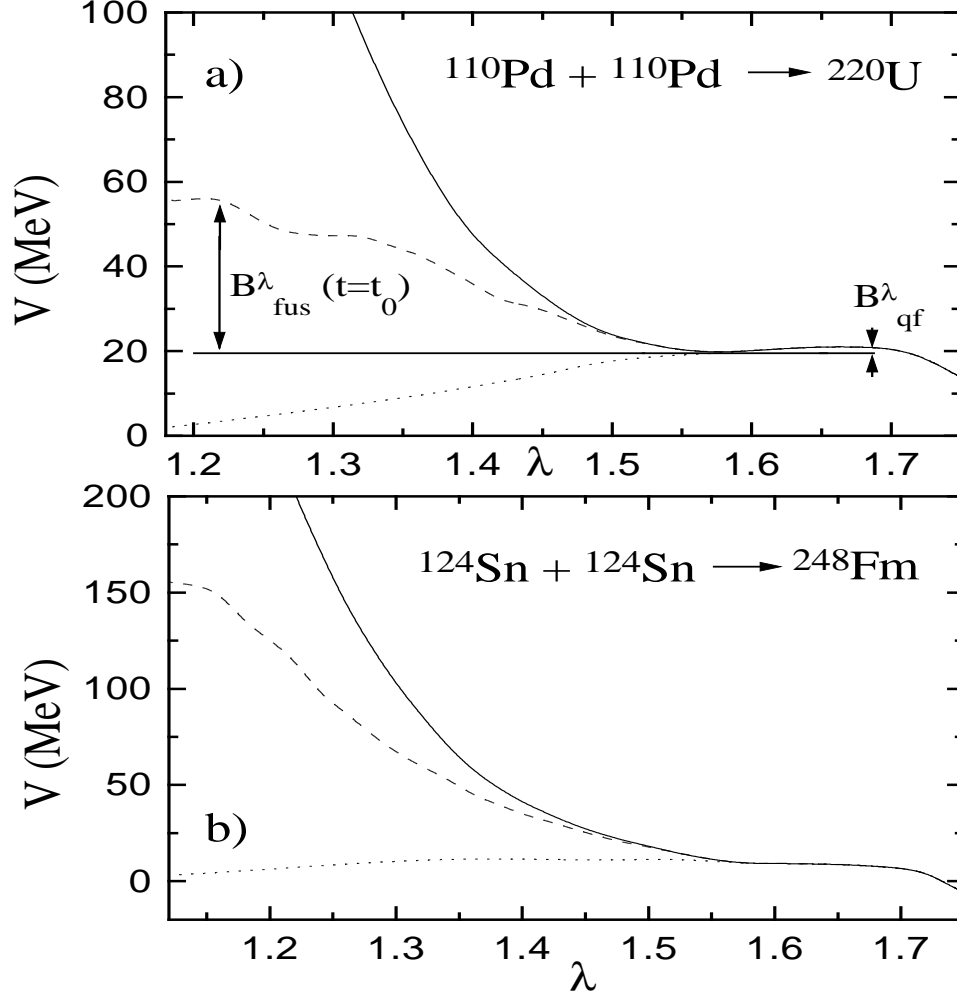| Reactions	1) P_{fus}^{λ}	2) P_{fus}^{λ}	P_{fus}^{η}	$P_{fus}^{\exp}$
^{90}Zr+^{90}Zr $\rightarrow^{180}$Hg	$2\cdot10^{-4}$	$2\cdot10^{-2}$	$2\cdot10^{-1}$	$\sim10^{-1}$
^{100}Mo+^{100}Mo $\rightarrow^{200}$Po	$9\cdot10^{-6}$	$3\cdot10^{-3}$	$2\cdot10^{-2}$	$5\cdot10^{-2}$
^{110}Pd+^{110}Pd $\rightarrow^{220}$U	$7\cdot10^{-15}$	$4\cdot10^{-7}$	$3\cdot10^{-4}$	$\sim10^{-4}$
^{56}Cr+^{164}Er $\rightarrow^{220}$U	$1\cdot10^{-6}$	$2\cdot10^{-3}$	$6\cdot10^{-1}$	
^{76}Ge+^{170}Er $\rightarrow^{246}$Fm	$9\cdot10^{-22}$	$3\cdot10^{-12}$	$6\cdot10^{-4}$	$8\cdot10^{-4}$
^{86}Kr+^{160}Gd $\rightarrow^{246}$Fm	$4\cdot10^{-26}$	$2\cdot10^{-16}$	$7\cdot10^{-5}$	$5\cdot10^{-5}$

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FIG. 1. a) Time-dependent dynamical potential $V(\lambda, t)$ as a function of elongation λ for the system ^{110}Pd+^{110}Pd. The initial diabatic potential $V(\lambda, t = 0) = V_{di}(\lambda)$ and the adiabatic potential $V_{ad}(\lambda)$ are shown by solid and dotted curves, respectively. The diabatic potential $V(\lambda, t = t_0)$ at the lifetime t_0 of the DNS is presented by a dashed curve. The nuclei are considered spherical with the neck parameter $\varepsilon = 0.74$. The parameter $\Gamma_0^{-1} = 0.030$ MeV^{-1} is used in the calculation of the single–particle widths. The fusion $B_{fus}^{\lambda}(t = t_0)$ and quasifission B_{qf}^{λ} barriers in λ are indicated. These barriers are measured with respect to the minimum of the potential. b) The same as in a) but for the system ^{124}Sn+^{124}Sn.
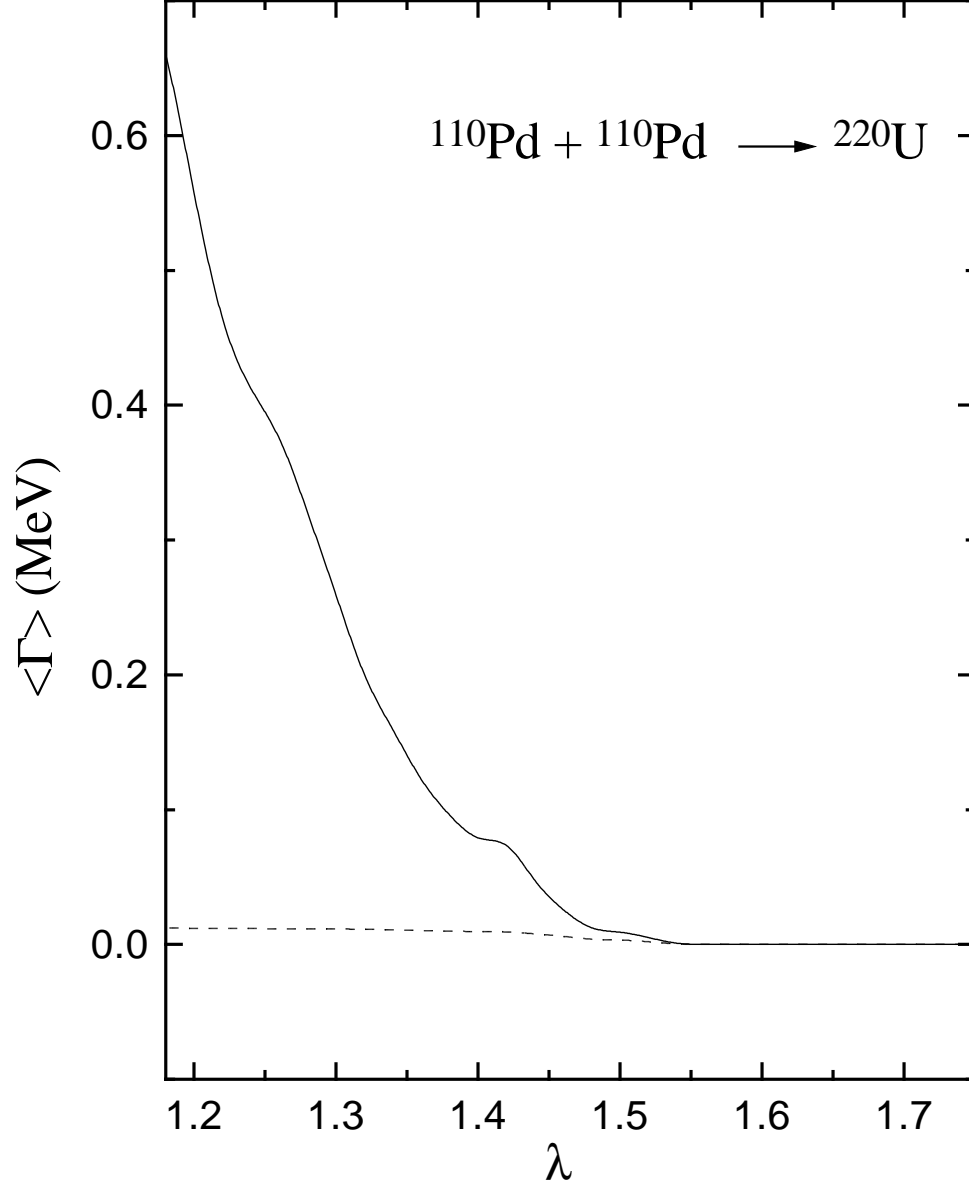
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FIG. 2. Time-dependent average width $< \Gamma(\lambda, t) >$ as a function of elongation λ for the system $^{110}Pd + ^{110}Pd$. The dependences $< \Gamma(\lambda, t = 0) >$ and $< \Gamma(\lambda, t = t_0) >$ are shown by solid and dashed curves, respectively.
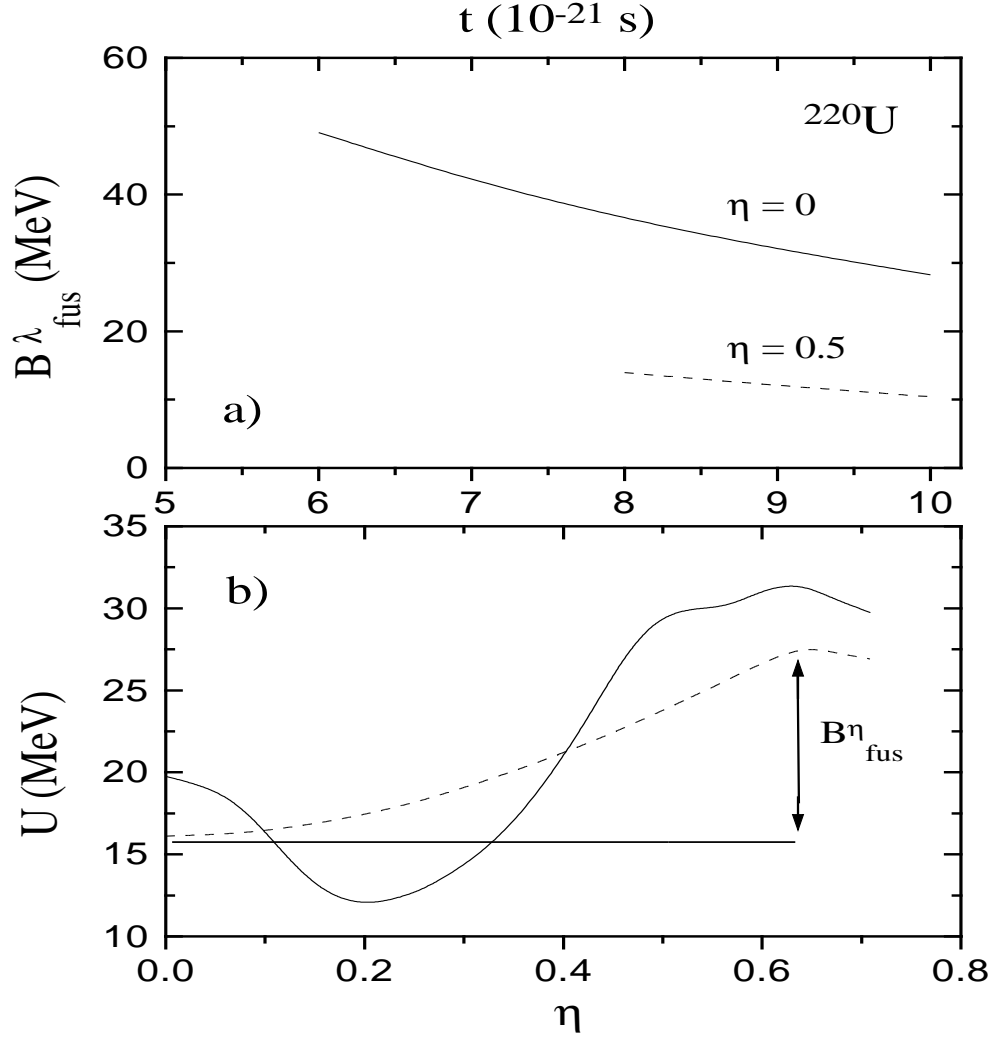
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FIG. 3. a) Inner fusion barriers $B_{fus}^{\lambda}(t)$ in λ as a function of time for the systems ^{110}Pd+^{110}Pd (solid curve) and ^{56}Cr+^{164}Er (dashed curve) which produce the same compound nucleus ^{220}U. The nuclei are considered spherical with $\varepsilon = 0.74$. The parameter $\Gamma_0^{-1} = 0.030$ MeV^{-1} is used in the calculation of single–particle widths. For times smaller than 6×10^{-21} s and 8×10^{-21} s for η=0 and 0.5, respectively, the potential $V(\lambda, t)$ is only repulsive and has no barrier (see Fig. 1). b) Calculated adiabatic potential energy of the DNS in the touching configuration of the nuclei as a function of mass asymmetry η for reactions leading to the same compound nucleus ^{220}U. The potential is calculated within the adiabatic TCSM with (solid curve) and without (dashed curve) shell corrections. The fusion barrier B_{fus}^{η} in η for the system ^{110}Pd+^{110}Pd is shown.